High Temperature Mixed State $c-$Axis Dissipation in Low Carrier Density $Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$

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The nature of the out-of-plane dissipation was investigated in underdoped $Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$ single crystals at temperatures close to the critical temperature. For this goal, temperature and angle dependent out-of-plane resistivity measurements were carried out both below and above the critical temperature. We found that the Ambegaokar-Halperin relationship [V. Ambegaokar, and B. I. Halperin, Phys. Rev. Lett. **22**, 1364 (1969)] depicts very well the angular magnetoresistivity in the investigated range of field and temperature. The main finding is that the in-plane phase fluctuations decouple the layers above the critical temperature and the charge transport is governed only by the quasiparticles. We also have calculated the interlayer Josephson critical current density, which was found to be much smaller than the one predicted by the theory of layered superconductors. This discrepancy could be a result of the d-wave symmetry of the order parameter and/or of the non BCS temperature dependence of the c-axis penetration length.

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Since the discovery of high temperature superconductivity in cuprates, the contrasting temperature T dependence of in-plane ρ_{ab} and out-of-plane ρ_c resistivities have been an issue of debate. This topic is even more complex in the underdoped systems where the density of states DOS of quasiparticles as well as the transfer integrals are momentum and temperature dependent. Additionally, strong fluctuations, which are expected at low carrier density, have a major contribution to dissipation.

The most debated issue is the interlayer dissipation and its field and angle dependence. In the most general way, the c-axis conductivity is presumably controlled by the tunneling of the Cooper pairs and quasiparticles, with conductivities σ_J and $\sigma_{c,qp}$, respectively; i.e., the c-axis conductivity $\sigma_c = \sigma_J + \sigma_{c,qp}$. This tunneling is the main consequence of the layered structure of the cuprates, which can be depicted as stacks of Josephson junctions made out of superconducting CuO_2 "electrodes" and intermediate blocking layers. Actually, this Josephson coupling of the layers distinguishes the layered cuprates from the ordinary anisotropic superconductors. The existence of this coupling was unambiguously demonstrated for cuprates with large anisotropy $\gamma > 100$, where γ is the ratio of the c-axis and inplane effective masses, $\gamma \equiv m_c/m_{ab}$, either directly by $I - V$ characteristics on small single crystals or messa structures[1,2,3,4,5,6] or by Josephson plasma resonance experiments.[7,8,9,10,11,12] In the case of systems with lower anisotropy, e.g., $YBa_2Cu_3O_{7-\delta}$ ($\gamma \approx 5-9$), the existence of a Josephson coupling in the c-axis direction was largely debated. Direct measurements were reported only in underdoped $YBa_2Cu_3O_{6+x}$.[13] while optical conductivity measurements[14,15] showed evidence of Josephson coupling in these low anisotropic cuprates.

As in the case of the in-plane transport, the interlayer dissipation is strongly influenced by phase fluctuations. A magnetic field applied perpendicular to the layers penetrates as pancake vortices with a particular phase distribution around each core. It is known that position fluctuations of the pancake vortices due to pinning and/or thermal diffusion of the vortex core, reduce phase correlations, hence, the Josephson coupling. Nevertheless, as long as phase correlations exist, they could provide a Josephson-type contribution to the out-of-plane transport. In fact, such phase correlations, though weak, were identified experimentally even in the liquid state of the vortex system.[8,16,17] Generally, these phase (vortex) fluctuations have important consequences both below and above the critical temperature T_{c0}. For $T < T_{c0}$, they drive the vortex system in a liquid state, whereas above T_{c0}, they allow vorticity to survive and to contribute to the in-plane dissipation.[18,19] To be specific, in the latter case, there is a crossover in the in-plane dissipation from a regime of pure flux-flow to a regime entirely due to quasiparticles, which occurs at a particular temperature *higher* than T_{c0}. The importance of the fluctuations increases when the density of charge carriers is reduced, i.e., in the case of underdoped cuprates.

Even though the interlayer coupling was investigated on a large extent in large-γ superconductors, the scarcity of data is evident in low and medium-anisotropic superconductors. Therefore, in the present study, using temperature, field, and angle dependence of the out-of-plane resistivity, we investigate the features of the c-axis dissipation in a medium-γ superconductor in an attempt to find the extent of Josephson response close to T_{c0}, i.e., in a temperature range where the phase fluctuations are important enough to reduce and/or sup-

press the Josephson coupling. As in the case of in-plane resistivity,[18,19] we take advantage of the different angular dependences of the different contributions to resistivity to obtain the desired information. This investigation is performed on $Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$ single crystals, where we used the antidoping effect of praseodymium to reduce the charge carrier density and increase the electronic anisotropy, hence, to drive the system in the strong fluctuation regime. In this way the electromagnetic and Josephson coupling of the pancake vortices and, subsequently, the interlayer coherence are weakened compared with the thermal fluctuations. Our main finding is that the c-axis dissipation scales with $Hcos\theta$ below T_{c0}, in a temperature range that ends at the critical temperature. We relate the failing of the scaling above T_{c0} to the suppression of the interlayer Josephson coupling by the in-plane phase fluctuations. So, even though in-plane strong superconducting fluctuations, hence, a large amount of condensate, persist above T_{c0},[18,19] there is no out-of-plane Josephson coupling which would provide an enhanced conductivity above T_{c0}.

$Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$ single crystals with typical dimensions of $1.0 \times 0.5 \times 0.02$ mm^3 were grown using a standard procedure described elsewhere.[20] The dimensions the the crystal for which the data are shown here are $0.6 \times 0.65 \times 0.017$ mm^3. We attached four gold wires (0.025 mm in diameter) with silver epoxy onto each of the two large faces of the single crystal (see top Inset to Fig. 1). The two outer (inner) contacts on the same face were used as current (voltage) terminals. The contact resistance is 2 Ω at room temperature. First, a magnetic field H up to 14 T was applied along the $c-$direction of the sample, a constant current $I \leq 1$ mA was fed through pads of size 0.05×0.5 mm^2, alternately on both faces, and the voltage on each face of the single crystal was measured at set temperatures between 0 and 300 K. Next, we measured in the same way the voltages at different constant temperatures but this time the single crystal was rotated in the applied magnetic field with the angle θ between H and the c-axis varying between 0 and 360°. The out-of-plane ρ_c and in-plane ρ_{ab} resistivities were calculated using an algorithm described elsewhere.[21] The critical temperature T_{c0} was taken at the midpoint of the normal-superconductor transition.

Figure 1 shows the temperature T dependence of ρ_c of an $Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$ single crystal measured at different applied magnetic fields, while its bottom Inset shows the zero field $\rho_c(T)$ over the whole measured temperature range. Even though this is a strongly underdoped sample, with a zero-field superconducting transition temperature $T_{c0} = 38$ K, it has a sharp transition, attesting to the good quality of this single crystal. The magnetoresistivity in the normal state is very small and positive. The normal state is metallic at high temperatures and becomes nonmetallic for temperatures lower than 133 K. The upturn in $\rho_c(T)$ observed below this temperature is the result of a complex process. First, there is a reduction in the planar density of states DOS due to the opening of the pseudogap at $\mathbf{k} = (\pi/2a, 0)$ (nodes) with decreasing T;[22] second, the transfer integral of the coherent contribution is angle dependent with maxima at the nodal points.

The angular dependence of the normalized $c-$axis resistivity $\rho_c(\theta)/\rho_c(\theta = 0°)$ of $Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$ measured at 30, 35 and 40 K in a magnetic field of 14 T is shown in Fig. 2. Both below and above T_{c0}, $\rho_c(\theta)$ displays a minimum at $\theta = 90°$ (i. e. for $H \parallel ab-$plane) and a maximum value at $\theta = 0°$ (i. e. for $H \parallel c-$axis). The former (latter) value of the angle corresponds to maximum (zero) transverse Lorentz force on the flux vortices. This fact rules out the possibility that flux motion contributes to the measured $c-$axis dissipation, since the measured dissipation is maximum for angles for which the Lorentz force is zero. Therefore, we assume that the c-axis transport in the mixed state involves Josephson and quasiparticle contributions which depend on T, H, and θ.

In the geometry we have used, the c-axis component of the current density flows mainly along the crystal edges in a wall of width almost equal to the pad width $\Delta = 50$ μm. We call this region the active area. The variation of the current density over this width is $[J_z(z, L/2) - J_z(z, L/2 - \Delta)]/J_z(z, L/2) < 14$ %, while, outside of this width, the current density decreases fast toward the center of the crystal. Therefore, we conclude that only these edge walls of the single crystal are involved in the current transport along c-axis.

The characteristic lengths for a stack of Josephson junctions are the Josephson penetration length λ_j, which accounts for the field penetration within the nonsuperconducting interlayer space and the c-axis London penetration length λ_c, which accounts for the magnetic screening. For this single crystal, $\lambda_j \approx 0.03$ μm, much smaller than Δ. Therefore, it is inappropriate to consider the single crystal as a stack of short junctions. For a stack of long junctions, Josephson vortices might be generated even in the absence of in-plane external fields if the in-plane currents are strong enough to create the required phase gradient. This phase gradient is generated if the junction length is larger than λ_c. Hence, if the width of the active area Δ is of the order of the c-axis magnetic penetration length λ_c one can assume that the in-plane currents are too small to generate Josephson vortices in the active area. There are no available data concerning λ_c for strongly underdoped $Y_{1-x}Pr_xBa_2Cu_3O_{7-\delta}$. An estimate of λ_c based on a similar underdoped $YBa_2Cu_3O_x$ is of the order of 30 μm.[23] Therefore, this Josephson system satisfies the condition $\lambda_j \ll \Delta \sim \lambda_c$. Hence, we assume that in this case the Josephson vortices are most likely absent and we can use the Ambegaokar-Halperin relationship. Additionally, the current we use is very low with a total current density $J = 4$ Acm$^{-2} \ll J_c$. Under these assumptions, we proceed to obtain an analytical relationship for the c-axis resistivity as a function of field, temperature, and angle.

Generally, the total conductivity can be derived from

Kubo's relationship for both Josephson and quasiparticle current.[24] However, this relationship is difficult to handle in the absence of an analytical dependence of the in-plane diffusion coefficient on field and temperature. Several experimental reports[25,26,27,28] have shown that the Ambegaokar-Halperin (AH) expression,[29] which is valid for a single Josephson junction, can be successfully used to fit the c-axis resistivity data when making specific assumptions on the expression of the critical current. The Ambegaokar-Halperin expression for c-axis resistivity is given by:

$$\rho_c(T) = \rho_n \left[\mathcal{I}_0\left(\frac{\Phi_0 I_c(T)}{2\pi k_B T}\right)\right]^{-2}, \quad (1)$$

where ρ_n is the intrinsic normal-state resistivity of the junction, $\mathcal{I}_0$ is the modified Bessel function, Φ_0 is the flux quantum, I_c is the critical current at a temperature T, and k_B is Boltzmann's constant. Note that the AH relationship accounts also for the contribution of the quasiparticles through ρ_n.

Because of the high energy of the Josephson coupling $\Phi_0 I_c/2\pi$ relative to the thermal energy, one can use the asymptotic expansion $\mathcal{I}_0(x) \approx \exp(x)/\sqrt{2\pi x}$ for $x > 1$. This approximation gives, for example, for $x = 3$ a 5% error compared with the exact Bessel function. An estimate of the zero-field conductivity of our samples gives $x(B = 0) = 5.5$ at 35 K and an error of approximately 2%. With this approximation, Eq. (1) gives the the following expression for the $c-$axis resistivity at high temperatures:

$$\rho_c(T, H) \approx \rho_{c,qp} \frac{\Phi_0 I_c(T, H)}{k_B T} \exp\left(-\frac{\Phi_0 I_c(T, H)}{\pi k_B T}\right). \quad (2)$$

Next, we discuss the T and H dependence of I_c. The temperature dependence of the AH relationship is limited only to the spin-wave type fluctuation of the order parameter. Therefore, we have to include in the above critical current term the contribution accounting for the presence of vortices. The maximum Josephson current I_c, which is related to the interlayer phase difference, is strongly influenced by the level of the fluctuations of the phases in each layer. To be specific, the Josephson current density J_c decreases as a result of both the thermally-induced misalignment of the planar vortices, which creates a gauge invariant phase difference $\varphi_{n,n+1}$ from layer to layer, and the thermally induced phase slippage (spin-wave type phase fluctuations). Regarding the thermal motion of the pancake vortices, there is a complex process of renormalization of J_c, which suppresses J_c.[30] A suppressed J_c increases in turn the penetration depth λ_c, hence, reduces the elastic constants, which in turn further suppresses J_c. This process is present both in the solid and liquid phases of the vortex system because the only difference between these two phases is the vanishing of the shear constant c_{66} in the latter phase. Additionally, the phase difference depends on the pancake position within each plane. Therefore, to obtain the interlayer critical current, one has to go beyond the ensemble average[31,32] and to also use a space average of the critical current. Hence,[33,34]

$$I_c^2 = J_{c0}^2 \int d\mathbf{r}_1 \int d\mathbf{r}_2 \exp\left\{i\left[\phi_{n,n+1}(\mathbf{r}_1) - \phi_{n,n+1}(\mathbf{r}_2)\right]\right\}, \quad (3)$$

with J_{c0} the local intrinsic (bare) Josephson critical current density. At high temperatures, Eq. (3) becomes:

$$I_c^2(T, H) = J_{c0}^2(T) A S(H),$$

where A is the area of the junction and $S(H)$, given by

$$S(H) = \int d\mathbf{r} \langle cos\left[\phi_{n,n+1}(\mathbf{r})\right] - cos\left[\phi_{n,n+1}(0)\right]\rangle, \quad (4)$$

is the correlation area, which needs to be evaluated. Following Koshelev, Bulaevski, and Maley,[31] we make the approximation

$$S(H) \approx f\gamma^2 s^2 \left(H_j/H_z\right)^{\alpha}. \quad (5)$$

Here, $f(H, T)$ is a function of order unity with a weak T dependence, s is the interlayer spacing, $H_j = \Phi_0/\gamma^2 s^2$ is a characteristic field, and $H_z = H cos\theta$ is the magnetic field in the c-axis direction. The deviation of the exponent $\alpha = 1 - k_B T/2\pi E_0(T)$ from unity is a result of the spin-wave type phase fluctuations and increases strongly close to the critical temperature. $[E_0 = s\Phi_0^2/(4\pi\mu_0\lambda_{ab}^2)$ is the Josephson energy of the area $\gamma^2 s^2]$.[35] The above approximation [Eq. (5)] is valid for applied magnetic fields larger than the characteristic field H_j. For the single crystals with a Pr doping $x = 0.46$, for which $\gamma = 26$ and $s = 11.7\mathring{A}$, the characteristic field $H_j \approx 2.3$ T. Hence, Eq. (5) is valid for $H > 2.3$ T. With this approximation, the Josephson critical current becomes

$$I_c(H, T) = J_{c0}(T)\gamma s f^{1/2} A^{1/2} (H_j/H_z)^{\alpha/2}. \quad (6)$$

Equation (4), hence Eq. (6), was derived in the approximation of a completely decoupled pancake vortex liquid, when the correlation of the $cos\varphi_{n,n+1}$ terms drops on a length scale of the order of the intervortex spacing.

Equations (2) and (6) give the following expression for the out-of-plane resistivity:

$$\rho_c(T, H, \theta) \approx \rho_{c,qp} \frac{\Phi_0 J_{c0}(T)\gamma s f^{1/2} A^{1/2}}{k_B T} \left(\frac{H_j}{H|cos\theta|}\right)^{\nu} \times$$

$$\exp\left[-\frac{\Phi_0 J_{c0}(T)\gamma s f^{1/2} A^{1/2}}{\pi k_B T} \left(\frac{H_j}{H|cos\theta|}\right)^{\nu}\right], \quad (7)$$

with $\nu = \alpha/2$. An important result of Eq. (7) is that the out-of-plane resistivity in the mixed state where Josephson tunneling dominates the conduction should scale with

$H\cos\theta$. Figure 3(a) is a plot of ρ_c vs $Hcos\theta$, measured at several temperatures below T_{c0} and in applied magnetic fields of 6, 8, 10, 12, and 14 T. Note that the data, indeed, follow the $H\cos\theta$ scaling.

Above the critical temperature, ρ_c vs $Hcos\theta$ plots do not map anymore onto a single curve [see Inset to Fig. 3(a) for $T = 40$ K]. This fact hints to the complete vanishing of the interplane phase correlations, hence, of the Josephson contribution to dissipation above T_{c0}. Thus, the phase fluctuations suppress the mechanism responsible for bulk superconductivity at the critical temperature, even though the in-plane dissipative processes still carry the hallmark of superconducting phase fluctuations up to temperatures well above T_{c0}.[18,19]

The term $\rho_{c,qp}$ in Eq. (7) ensues from the quasiparticle current driven by the time variation of the gauge invariant phase difference. In the case of d-wave superconductors, the quasiparticle concentration does not vanish with decreasing temperature. At any temperature T, there is always a $\mathbf{k}$ range so that $\Delta(\mathbf{k}) < k_BT$, which facilitates the quasiparticle excitation near the gap nodes. Microscopic models have shown that in the case of constant DOS, the quasiparticle out-of-plane resistivity below the critical temperature depends on temperature as $\rho_{c,qp} = \rho_n(3\Delta_0^2/\pi T^2)$ if the tunneling is coherent and is T-independent if the tunneling is completely incoherent.[36] A real material displays both contributions, hence, it follows a power law temperature dependence. Additionally, the DOS decreases with decreasing T in underdoped cuprates. The absence of an analytical expression for the temperature dependence of the DOS makes impossible the determination of the temperature dependence of the quasiparticle contribution to resistivity. In the presence of a magnetic field, there is a small change in conductivity due to the Doppler shift in the quasiparticle spectrum.[37] However, a sensitive change requires extremely high magnetic fields,[38] so that in the present measurements ($H \leq 14$ T) $\rho_{c,qp}$ is practically field independent. Indeed, the c-axis resistivity data show that the normal-state magnetoresistivity is very small, which implies an almost H independent quasiparticle contribution to the out-of-plane conduction. Therefore, we assume that $\rho_c(T)$ measured in 14 T in the normal state and its extrapolation at lower temperatures in the mixed state is the out-of-plane quasiparticle resistivity $\rho_{c,qp}(T)$.

With $\rho_{c,qp}(T)$ determined as just discussed above, we fit the $\rho_c(T,H,\theta)$ data with Eq. (7) with two fitting parameters: the exponent $\nu(T)$ and

$$C(T) = \frac{\Phi_0 J_{c0}(T)\gamma s A^{1/2} H_j^{\nu(T)}}{\pi k_B T},$$

where we take $f \approx 1$. The results of the fitting of the data measured at several temperatures, and in 14 T and 10 T are shown in Fig. 3(b) and its Inset, respectively. The excellent fit of the out-of-plane resistivity data with Eq. (7) confirms the validity of our approach and shows that the T, H, and θ dependence of the measured out-of-plane resistivity in the mixed state is dominated by the Josephson tunneling of the Cooper pairs and the quasiparticle tunneling.

The picture that immerges from these results is as follows. As in the case of the in-plane dissipation, the fluctuations have a significant effect on the nature of ρ_c at high temperatures in low charge carrier density cuprates such as $Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$. Nevertheless, although the dissipations along the two directions have the same origin, they are governed by different mechanisms, hence, display two temperature scales. The out-of-plane dissipation is governed by the Josephson tunneling of the Cooper pairs and the quasiparticle tunneling. With increasing T, the in-plane phase fluctuations give rise to a rapid suppression of the Josephson coupling, the unique process which makes superconductivity a bulk phenomenon, and of the corresponding interlayer supercurrent density. Small interlayer correlations survive above the irreversibility temperature up to 37 K above which both the Josephson coupling and the corresponding interlayer supercurrent vanish. At $T > T_{c0}$, even though in-plane vorticity has been shown to exist in a strong fluctuating regime,[18] these phase fluctuations are too fast to allow any phase correlation along the c-axis. Hence, $\sigma_J = 0$ and the out-of-plane conductivity is only a result of quasiparticle tunneling. This is not the case of the in-plane dissipation, which displays a contribution from phase fluctuations (pancake vortices) arising from their motion driven by the transport current up to a charge carrier density dependent temperature $T_\varphi >> T_{c0}$.[18,19] Therefore, the temperature scale is T_{c0} (T_φ) for the contribution of the superconducting dissipation to the total out-of-plane (in-plane) resistivity. These two temperatures merge as the density of charge carriers increases toward the optimal doping.

Additional improvements to the model used in this study would require the incorporation of nonequilibrium effects due to the nodal quasiparticles, mainly in the high temperature range.[39] This simple AH approach, which constitutes the starting point of our analysis, though fruitful, has been the subject of criticisms regarding the omission of the interaction between adjacent junctions,[40] and the assumption of a Fermi liquid behavior in underdoped cuprates. However, the simplicity of the AH relationship and its reported success at low temperatures make it attractive, with appropriate assumptions, in the high temperature regime. Our present results confirm, indeed, its applicability close to T_{c0}.

From the angular magnetoresistivity data, we also extracted the temperature dependence of the bare interlayer (Josephson) critical current density at high temperatures, close to T_{c0}, from

$$J_{c0}(T) = \frac{\pi k_B T C(T)}{\Phi_0 \gamma s A^{1/2} H_j^{\nu(T)}}, \tag{8}$$

using the fitting parameter $C(T)$ as obtained from the fit of the angular magnetoresistivity data at high temperatures with Eq. (7). A plot of J_{c0} vs T is shown in Fig. 4. These values of J_{c0} are smaller than the values smaller

than the values predicted by a simple model of layered superconductors, which gives $J_{c0} = \Phi_0/[2\pi\mu_0\lambda_c^2(T)]$. The temperature dependence of the data follow the power law $J_{c0}(T) \approx 7.6(T/T_{c0})^{-1.73}$. Such a T dependence could be the result of the complexity of the interlayer Cooper pair transport in cuprates containing conducting CuO chains combined with the $d-$wave symmetry of the superconducting order parameter. Therefore, the temperature dependence of J_{c0} is provided not only by $\lambda_c^{-2}(T)$, but also by the T-dependence of the density of states of the localized resonant centers.[41] The $\lambda_c^{-2}(T)$ itself changes its convexity at high temperatures,[23] most probably due to the excitation of the quasiparticles out of the condensate at gap nodes.

The temperature dependence of the exponent ν is shown in the Inset to Fig. 4. A fit of the data gives a linear T dependence, i.e., $\nu(T) = 1.7(1 - T/T^*)$, where $T^* = 39.5$ K is slightly higher than T_{c0} defined as the mid point of the transition curve. Theoretically, the exponent ν should be linear in $T\lambda_{ab}^2(T)$. A plot of the theoretical $\nu(T)$ curve using the BCS dependence of $\lambda_{ab}^2(T)$ is also shown in Fig. 4. The experimental and theoretical values are close to each other for $T \geq 30$ K. At lower temperatures, the data obtained from fitting are almost twice as high as the theoretical values. Actually, the expression used for the field dependence of I_c [Eq. (6)] is not valid at low temperatures. This could be one reason for the above discrepancy. It is interesting to note, however, that the scaling of $\rho_c(H|cos\theta|)$ still works down to 20 K. Another reason for the discrepancy between the experimental and theoretical values of ν could be that $\lambda_{ab}^{-2}(T)$ has a non BCS temperature dependence due to the nodal quasiparticles.

In summary, we analyzed the out-of-plane dissipation in a medium anisotropic underdoped cuprate at temperatures around T_{c0}. We performed these measurements in order to investigate the origin of the large ρ_c and its T, H, and angle dependence in this material. The data are well fitted by the Ambegaokar-Halperin expression for temperatures up to the critical temperature and applied magnetic fields as high as 14 T. We found that the interlayer resistivity follows a simple scaling law as a function of magnetic field and angle; i.e., $\rho_c(H, \theta) = \rho_c(|Hcos\theta|)$. The existence of the scaling close to the critical temperature proves the persistence of interlayer correlations above the irreversibility temperature. Nevertheless, the scaling fails above the mid point critical temperature, above which the c-axis charge transport is governed by quasiparticles only. This is different from the in-plane dissipation, in which the contribution of the superconducting fluctuations can be discerned up to temperatures as high as $1.5 \times T_{c0}$. We also have determined the interlayer critical current density. It was found to be lower than predicted by simple models of Josephson coupled superconductors.

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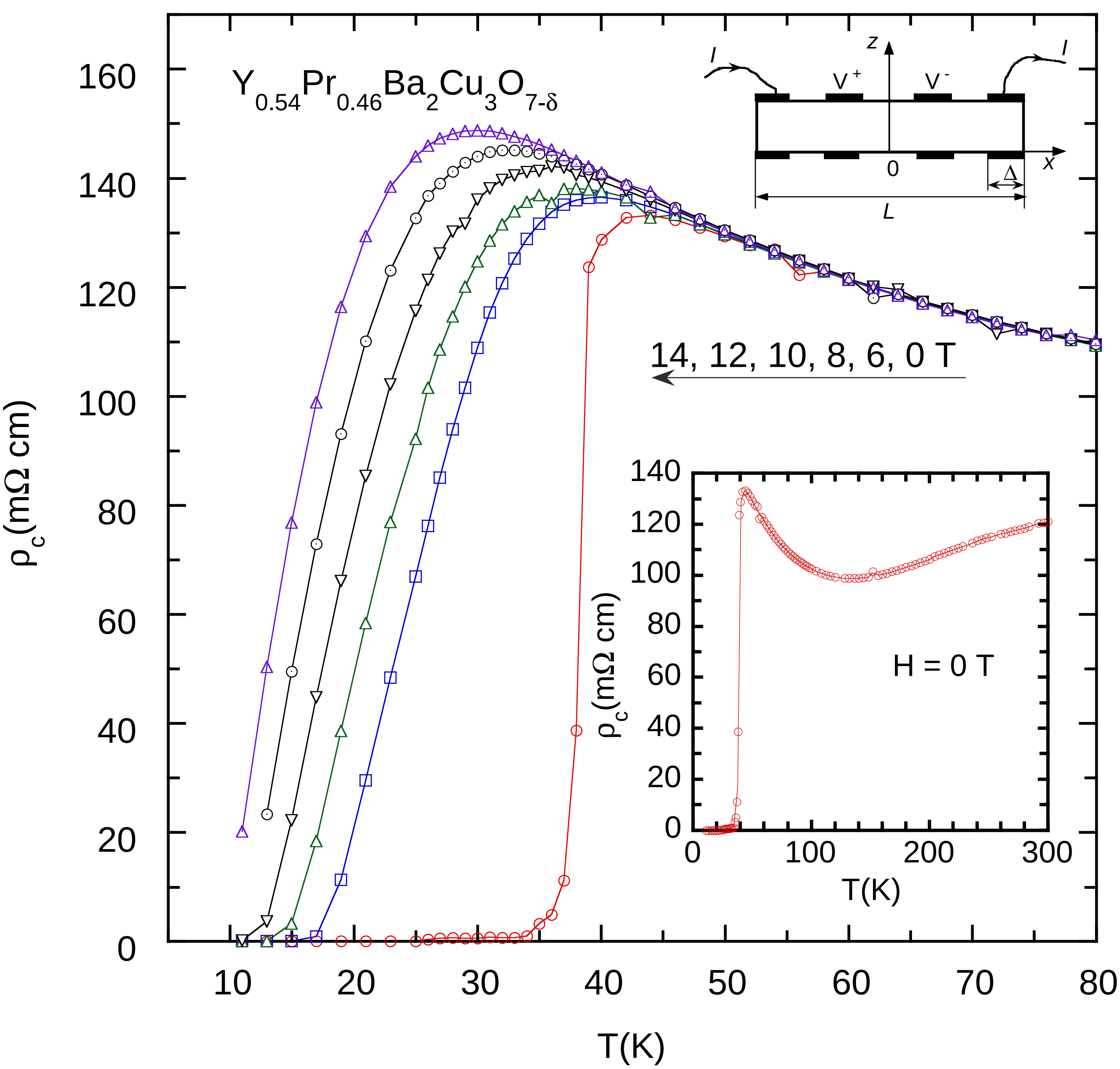
FIG. 1: Temperature T dependence of out-of-plane resistivity ρ_c of an $Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$ single crystal, measured in an applied magnetic field of 0, 6, 8, 10, 12, and 14 T and for $T \leq 80$ K. Insets: (top) Sketch of sample geometry and leads configuration. (bottom) Zero field $\rho_c(T)$ shown over the whole measured T range. The solid lines are guides to the eye.

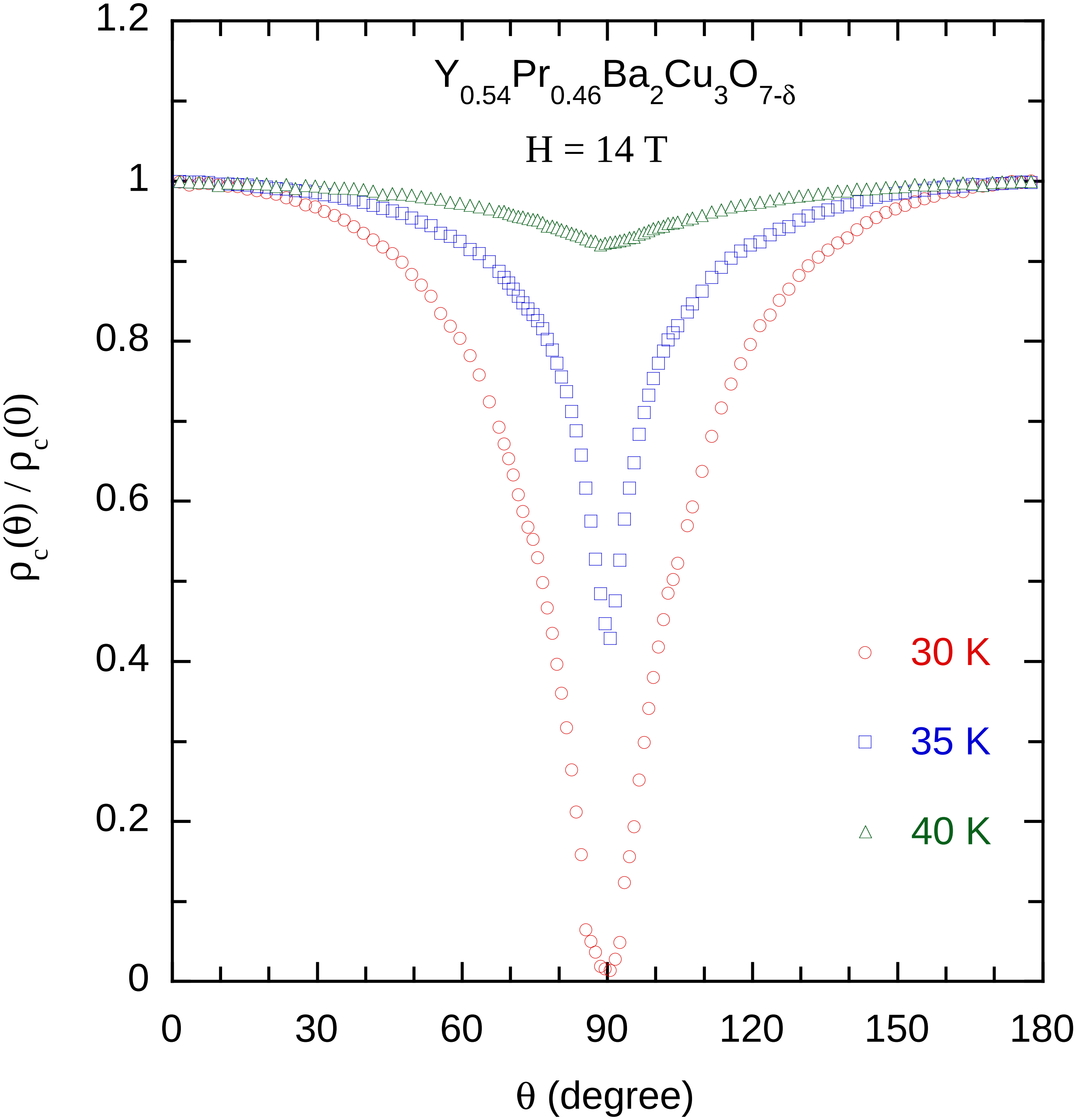
FIG. 2: Angular θ dependence of normalized out-of-plane resistivity $\rho(\theta)/\rho(0)$ of an $Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$ single crystal, measured at 30, 35, and 40 K and 14 T.

FIG. 3: (a) Plot of the out-of-plane resistivity ρ_c vs $H|cos\theta|$ of an $Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$ single crystal, measured at 25, 30, 35, and 36 K and 6 T (open circles), 8 T (diamonds), 10 T (triangles), 12 T (inverted triangles), and 14 T (open squares). Inset: Plot of $\rho_c(H|\cos\theta|)$ measured at 40 K. (b) and its Inset: Same plot of the data measured at 14 T and 10 T, respectively. The solid lines are fits of the data with Eq. (7).

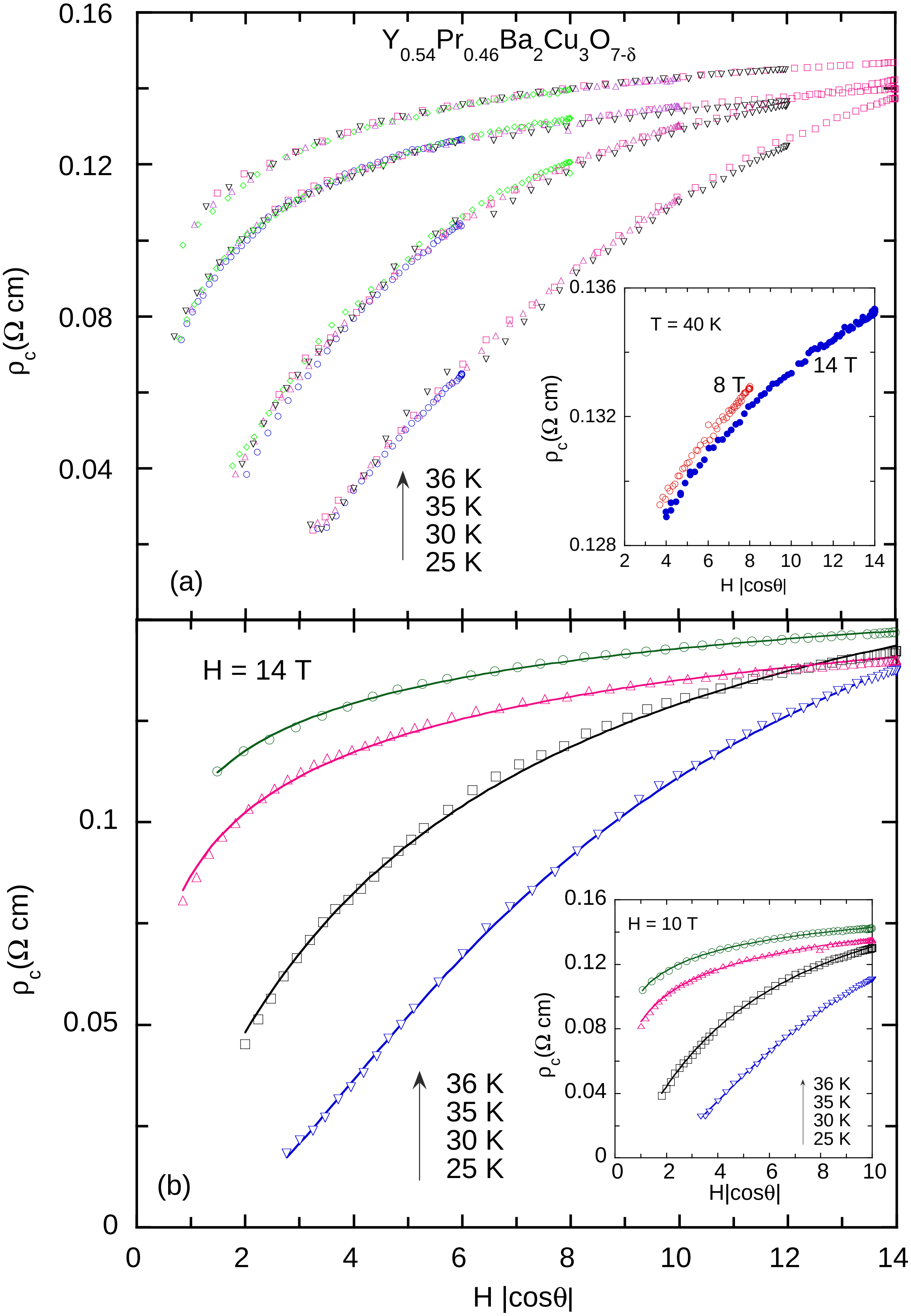
FIG. 4: Josephson critical current density J_{c0}, calculated with the fitting parameters obtained by fitting the angular magnetoresistivity, vs reduced temperature T/T_{c0}. The solid line is a power-law fit. Inset: Temperature T dependence of exponent ν (empty circles) obtained by fitting the magnetoresistivity data. The theoretical $\nu(T)$ dependence is calculated with the T dependence of the penetration depth $\lambda_{ab}(T)$ given by the BCS theory and taking as the critical temperature $T_{c0} = 38$ K.



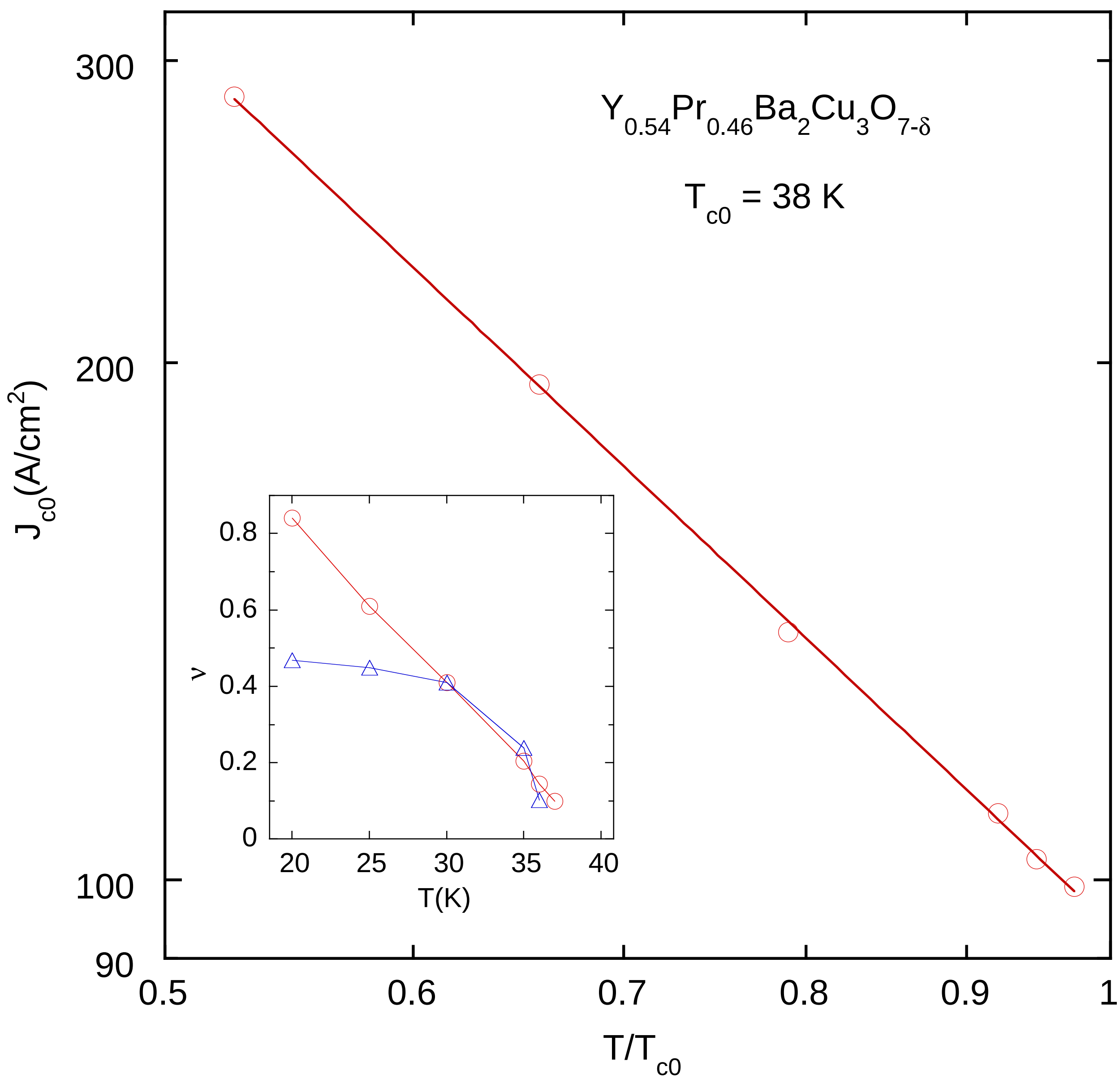
$Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$
14, 12, 10, 8, 6, 0 T
ρ_c(mΩ cm)
T(K)
H = 0 T
I
V+
V-
z
x
0
Δ
L



$Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$
H = 14 T
30 K
35 K
40 K
$\rho_c(\theta) / \rho_c(0)$
θ (degree)
0
0.2
0.4
0.6
0.8
1
1.2
0
30
60
90
120
150
180



$Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$
$\rho_c(\Omega\ cm)$
0.16
0.12
0.08
0.04
T = 40 K
8 T
14 T
0.136
0.132
0.128
2 4 6 8 10 12 14
H |cosθ|
36 K
35 K
30 K
25 K
(a)
H = 14 T
0.1
0.05
0
H = 10 T
0.16 0.12 0.08 0.04 0
0 2 4 6 8 10
H|cosθ|
(b)
0 2 4 6 8 10 12 14
H |cosθ|



$Y_{0.54}Pr_{0.46}Ba_2Cu_3O_{7-\delta}$
T_{c0} = 38 K
$J_{c0}(A/cm^2)$
T/T_{c0}
300
200
100
90
0.5
0.6
0.7
0.8
0.9
1
ν
T(K)
0.8
0.6
0.4
0.2
0
20
25
30
35
40